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Baryogenesis in Brans-Dicke theory

Sietse van der Post* and Tomislav Prokopec†

Institute for Theoretical Physics (ITP) & Spinoza Institute, Utrecht University,

Leuvenlaan 4, Postbus 80.195, 3508 TD Utrecht, The Netherlands

Abstract

A new mechanism for baryogenesis is proposed in the context of an extended Brans-Dicke (BD) theory. We generalize the BD scalar to complex field with CP violating coupling to curvature and show that the charged BD current can be enhanced during inflationary epoch. After inflation the current decays (via tree level interactions) into the standard model particles. When the BD scalar is charged under baryon and/or lepton number, the decay produces a net baryon number. Rather generically a sufficiently large scalar current can be produced during inflation to account for the observed baryon asymmetry of the Universe. Our baryogensis scenario can in an elegant way be incorporated into a model of extended inflation.

*post@phys.uu.nl

†T.Prokopec@phys.uu.nl

I. INTRODUCTION

In an expanding universe particle densities scale rather generically as $n \propto 1/a^3$, where a denotes the scale factor of the Universe. This observation immediately implies that any charge density preceding the inflationary epoch is diluted during inflation into oblivion. Namely the volume during inflation expands so tremendously that information about any initial charge is forgotten during inflation. This is true of course only if the dilution is not compensated by the production. To our knowledge up to now no mechanism has been proposed by which a charge density in some particles would grow during inflation. Here we propose such a model within the context of an extended Jordan-Fierz-Brans-Dicke (BD) theory [1], where the BD scalar is generalised to a complex scalar field. We shall argue that under rather generic conditions, an initial small scalar charge density extending over several Hubble volumes can be efficiently amplified during inflation. At the end of inflation the charge decays into ordinary matter fields, converting thus the scalar charge into baryonic or/and leptonic charges, which is in turn reprocessed by sphalerons into a baryon asymmetry, representing thus a mechanism for creation of the observed baryon asymmetry of the Universe.

The use of scalar fields has been ubiquitous in baryogenesis models, ranging from Affleck-Dine baryogenesis [2, 3], spontaneous baryogenesis [4], to electroweak scale baryogenesis [5, 6]. Yet with a few exceptions [7, 8], inflation has not been used in connection with baryogenesis.

The attractive feature of our baryogenesis model is that the field responsible for the strength of the gravitational coupling is also responsible for the generation of matter. In General Relativity *no* such mechanism is possible since the metric tensor cannot violate CP symmetry.

The plan of the Letter is simple. In section II we define our generalized Brans-Dicke theory. The resulting equations are solved and analysed in section III and in section IV we conclude.

II. BRANS-DICKE THEORY WITH A COMPLEX SCALAR

In Brans-Dicke theory the gravitational constant G is promoted to a dynamical parameter changing in space and time [1]. This variation is conventionally expressed through a scalar field Φ such that $G = G_0\Phi^{-1}$ with G_0 the bare, true constant of nature. The Brans-Dicke action is in natural units ($c = G_0 = \hbar = 1$) given by,

$$S = \frac{1}{16\pi}\int d^4x(-g)^{1/2}\left[\Phi R - \frac{\omega}{\Phi}(\partial_\mu\Phi\partial_\nu\Phi)g^{\mu\nu}\right] + \int d^4x\mathcal{L}_{\rm mat}\,. \qquad (1)$$

Here R denote the Ricci scalar curvature, $\mathcal{L}_{\rm mat}$ is the ordinary Lagrangian of matter and non-gravitational fields, ω is the BD coupling constant and $g = \det(g_{\mu\nu})$. For dimensional reasons as well as for convenience the substitution $\Phi \to \phi\phi$ is sometimes made in literature. We will make an analogous substitution, since it converts the kinetic term to the canonical form.

Here we propose the following generalisation of the Brans-Dicke theory. We promote ϕ to complex scalar and write the scalar sector of the theory (1) as follows,

$$(-g)^{-1/2}\mathcal{L}_\phi = -\rho(\partial_\mu\phi\partial_\nu\phi^*)g^{\mu\nu} + \Big(-\frac{1}{2}\bar{\omega}(\partial_\mu\phi\partial_\nu\phi)g^{\mu\nu} + \bar{\mu}\phi^2 R + \text{h.c.}\Big)\,. \tag{2}$$

Note that this lagrangian contains operators up to dimension four only, and in this sense the renormalisability properties of gravity are not made worse by the theory (2). Furthermore, it can accomodate CP violation when the (dimensionless) parameters $\bar{\omega}$ and $\bar{\mu}$ are complex,

$$\bar{\omega} = \omega \mathrm{e}^{i\theta_\omega}, \quad (\omega \equiv |\bar{\omega}|)\,, \qquad \bar{\mu} = |\bar{\mu}|\mathrm{e}^{i\theta_\mu}\,. \tag{3}$$

There is only one physical CP-violating phase in Eq. (2), however. Indeed, the phase of $\bar{\omega}$ can be removed by a field rotation, $\phi \to \phi\exp(-i\theta_\omega/2)$, upon which the kinetic terms in Eq. (2) become real. The remaining complex parameter,

$$\mu \equiv \bar{\mu}\mathrm{e}^{-i\theta_\omega} = \mu_r + i\mu_i\,, \tag{4}$$

whose components are of the form,

$$\mu_r = |\mu|\cos(\theta_\mu - \theta_\omega)\,, \qquad \mu_i = |\mu|\sin(\theta_\mu - \theta_\omega)\,, \tag{5}$$

mediates CP violation of the theory (2). Note in particular that, when $\theta_\mu - \theta_\omega \neq n\pi$ (n =integer), then $\mu_i \neq 0$, and there is a nonvanishing CP violation in the theory. Moreover, when $\pi/2 < \theta_\mu - \theta_\omega < 3\pi/2$, then $\mu_r < 0$. As we shall see, the case of negative μ_r's plays an important role in the analysis that follows.

Dubbing the real and imaginary component of the rotated field as $\phi_\pm$ such that,

$$\phi\mathrm{e}^{-i\theta_\omega/2} = \phi_+ + i\phi_-\,, \tag{6}$$

we can rewrite lagrangian (2) as,

$$(-g)^{-1/2}\mathcal{L}_\phi = -(\rho+\omega)(\partial_\mu\phi_+\partial_\nu\phi_+)g^{\mu\nu} - (\rho-\omega)(\partial_\mu\phi_-\partial_\nu\phi_-)g^{\mu\nu} + 2\mu_r(\phi_+^2 - \phi_-^2)R - 4\mu_i\phi_+\phi_- R\,. \tag{7}$$

The coefficients of the kinetic terms have to be such that $\rho - \omega > 0$ in order to have positive energy states. As can be seen the original lagrangian corresponds to two fields with different dynamics. The presence of charge violation should be clear by the mixing between the two fields in the last term and the sign difference in the quadratic terms. We emphasize that, although this model can be written in terms of two coupled real scalars, the physics is different. The two fields are connected through the relation $\phi e^{-i\theta_\omega/2} = \phi_+ + i\phi_-$, such that a non-vanishing scalar current can arise.

Since in our model baryon number is created via a postinflationary decay of a scalar current $j^0_\phi \equiv Q$, we shall mainly be interested in the production of a scalar charge density q_ϕ during inflation,

$$q_\phi \equiv J_\phi = \frac{i}{2}(\phi\dot{\phi}^* - \phi^*\dot{\phi}) \,. \tag{8}$$

A scalar field is said to be charged under the baryon number B (or any other charge not proportional to B+L) if – when it decays via tree level processes into the standard model matter fields – it generates a net lepton an baryon number according to the relation,

$$Q = q_\phi V = q_B B + q_L L \,, \tag{9}$$

where Q denotes the total scalar charge, V is the spatial volume, B and L denote the baryon and lepton number densities, respectively, and q_B and q_L are the corresponding charges. Sphalerons [9, 10] then process any $(B-L)_0$ charge produced during inflation into a net B and L according to [11],

$$\begin{aligned} B &= \frac{8n_f + 4(n_H+2)}{24n_f + 13(n_H+2)}(B-L)_0 \\ L &= -\frac{16n_f + 9(n_H+2)}{24n_f + 13(n_H+2)}(B-L)_0 \,, \end{aligned} \tag{10}$$

where n_F and n_H denote the number of quark and lepton families and Higgs doublets, respectively. This can be then observed as the baryon and lepton asymmetry of the Universe today.

III. ANALYSIS OF THE MODEL

We shall now analyse the dynamics of the scalar field governed by the lagrangian (2) (and (7)) during inflation. Inflation in a Brans-Dicke theory [12] is of a powerlaw type with the Ricci curvature scalar decaying as $R \propto t^{-2}$. The solutions for the scalar Φ and corresponding scale factor in the original BD theory of Eq. (1) are,

$$\begin{aligned} \Phi &= M_p^2(1+\chi t/\alpha)^2 \\ a(t) &= a_0(1+\chi t/\alpha)^{\omega+1/2} \,, \end{aligned} \tag{11}$$

where $M_p = (8\pi G)^{-1/2}$ denotes the Planck mass, $\chi^2 = 8\pi\rho_f/3M_p^2$ the Hubble constant squared in the Einstein frame, $\alpha^2 = (3+2\omega)(5+6\omega)$, and ρ_f is the energy density driving inflation. As can be seen, for early times and large ω the expansion is nearly exponential $a(t) \approx e^{\chi t}$ and R is decays slowly. For simplicity here we assume that ω in Eq. (11) is large, such that we can approximate the space-time by de Sitter space with the metric tensor given by

$$g_{\mu\nu} = \mathrm{diag}(-1, a^2, a^2, a^2)\,, \qquad a = a_0 \exp(Ht)\,, \tag{12}$$

where a denotes the scale factor (a_0 is some initial scale factor of no physical relevance) and H is the Hubble parameter.

In de Sitter space (12) the lagrangian (7) implies the following equation of motion for the homogeneous modes $\phi_\pm = \phi_\pm(t)$,

$$\ddot{\phi}_\pm + 3H\dot{\phi}_\pm \mp 2\xi_\pm R\phi_\pm + \zeta R\phi_\mp = 0\,, \tag{13}$$

where

$$\xi_\pm = \frac{\mu_r}{\rho \pm \omega}\,, \qquad \zeta = \frac{2\mu_i}{\sqrt{\rho^2 - \omega^2}}\,. \tag{14}$$

Note that we are considering the evolution of charged configurations which are to a good approximation homogeneous at the beginning of inflation. If the Universe was highly inhomogeneous at the beginning of inflation, only a small fraction of field configurations qualify, such that the corresponding initial amplitude may be small [7, 8]. The amplitude of these large scale charge fluctuations, which we assume to be of statistical origin, can be estimated as follows. The number of particles N and antiparticles $\bar{N}$ in a volume $V \sim H^{-3}$ are given by $N \sim T^3 V$, $\bar{N} \sim T^3 V$, and the corresponding charge is $Q \sim N - \bar{N}$, where T denotes the effective temperature (the system is not assumed to be in thermal equilibrium). To estimate the amplitude of charge fluctuations δQ, note that because of total charge neutrality of the system, each charge is paired with an opposite charge, such that the boundary surface $A = \partial V$ contains $N_A \sim A/\ell^2 \sim T^2 V^{2/3}$ charge pairs, where $\ell \sim T^{-1}$ is the average charge separation. The statistical charge fluctuation $Q_A = \sqrt{N_A}$ then corresponds to the expected net charge in the volume V, $\delta Q \sim Q_A \sim T V^{1/3}$. Given that $V^{1/3} \sim 1/H$ and $T \sim \sqrt{H M_P}$, we get for the charge fluctuation in a Hubble volume $\delta Q \sim (M_P/H)^{1/2}$, which gives for the initial current, $j_\phi \sim \delta Q/V \sim H^3 (M_P/H)^{1/2}$. The charge fluctuation δQ is of course suppressed when compared with the total particle number $N \sim (M_P/H)^{3/2}$ as $\delta Q/N = j_\phi/(N/V) \sim H/M_P$. For example, when this estimate is applied to the electroweak scale, one gets for the baryon to entropy ratio, $n_B/s \sim (T/M_P)^2 \sim H/M_P \sim 10^{-34}$. From this estimate and the analysis presented below it then follows that – as long as the charge does not decay during inflation – the preinflationary charge density fluctuations should suffice to make our baryogenesis mechanism effective.

In de Sitter space (12) Eq. (13) can be solved by an exponential *Ansatz*. Since the equations are coupled, there are four independent solutions with corresponding real constants $\phi^{(0)}_{\pm j}$. The solutions read,

$$\begin{aligned}
\phi_\pm &= \Re \sum_{j=1}^{4} \phi^{(0)}_{\pm j} e^{\kappa_j t}, \qquad \kappa_{\pm\pm} = -\frac{3}{2} H \left(1 \pm \sqrt{1-\beta\pm\gamma}\right) \\
\beta &= \frac{8R\omega\mu_r}{9H^2(\rho^2-\omega^2)}, \qquad \gamma = \frac{8R}{9H^2(\rho^2-\omega^2)}\sqrt{\rho^2\mu_r^2+(\rho^2-\omega^2)\mu_i^2}\,.
\end{aligned} \tag{15}$$

Note first that only four out of the eight constants $\phi^{(0)}_{\pm j}$ are independent. This is so because $\phi^{(0)}_{+j}$ and $\phi^{(0)}_{-j}$ are related by Eq. (13) as follows,

$$[\kappa_j(\kappa_j+3H) - 2\xi_+ R]\phi^{(0)}_{+j} + \zeta R \phi^{(0)}_{-j} = 0\,. \tag{16}$$

Secondly, as can easily be verified, only κ_{--} and κ_{-+} can be positive and hence can give growing solutions. For certain choice of parameters κ_{--} can become complex, though. In order to have a growing current, we need two independent solutions, since the contribution of solutions of $\phi_\pm$ with equal time behavior is zero as can be seen from Eq. (8). The leading contribution therefore comes from the interference between κ_{--} and κ_{-+}. Neglecting the other two decaying terms, we adopt the solution,

$$\phi_\pm = \Re\left[\phi^{(0)}_{\pm 1} e^{\kappa_{--}t} + \phi^{(0)}_{\pm 2} \mathrm{e}^{\kappa_{-+}t}\right], \tag{17}$$

where we took the real part since $\phi_\pm$ are by definition real. Substituting this in definition (8) of the current we arrive at,

$$J_\phi = (\phi^{(0)}_{+1}\phi^{(0)}_{-2} - \phi^{(0)}_{+2}\phi^{(0)}_{-1})\Re\left\{(\kappa_{-+}-\kappa_{--})\,\mathrm{e}^{(\kappa_{--}+\kappa_{-+})t}\right\}\,, \tag{18}$$

where we took account of the fact that in order to get a growing solution, $\Re[\kappa_{--}+\kappa_{-+}] > 0$, implying that κ_{-+} must be real.

There are two cases of interest:

`Case A.` When $1-\beta-\gamma<0$, then κ_{--} is complex and κ_{-+} is real with $1-\beta+\gamma>4$, which implies:

$$1-\gamma < \beta < \gamma - 3\,, \tag{19}$$

from which we also infer that $\gamma > 2$. In this case the current oscillates with a growing envelope amplitude. In order to prevent fast oscillations, some tuning is required. In particular, the oscillations are slow if $\beta \approx 1-\gamma$. In this case $\beta<0$ and $\mu_r<0$, since $\gamma>2$. The corresponding current is

$$J_\phi = J_0^{(1)}\left[\cos(i\alpha_- Ht) + \frac{i\alpha_-}{\alpha_+}\sin(i\alpha_- Ht)\right]\mathrm{e}^{(\alpha_+-3)Ht}\,, \tag{20}$$

where

$$\alpha_+ = \frac{3}{2}\sqrt{1-\beta+\gamma}\,, \qquad \alpha_- = \frac{3}{2}\sqrt{1-\beta-\gamma}\,, \qquad J_0^{(1)} = H\alpha_+(\phi_{+1}^{(0)}\phi_{-2}^{(0)} - \phi_{+2}^{(0)}\phi_{-1}^{(0)})\,. \qquad (21)$$

A typical evolution of the current (20–21) is shown in figure 1.

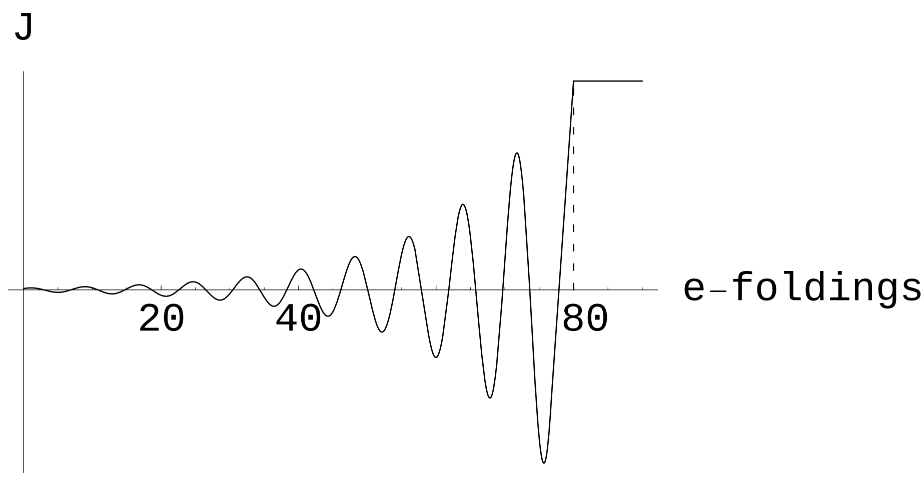

FIG. 1: The scalar charge density (current) during inflation as a function of the cosmic time (the number of e-foldings). Note oscillations with a growing amplitude. After inflation, the charge density freezes out.

`Case B.` When $1-\beta-\gamma > 0$, both κ_{--} and κ_{-+} are real. In this case the current grows provided the condition

$$\beta < -\frac{\gamma^2}{4} \qquad (22)$$

is fulfilled. In this case no oscillations are present. From Eq. (22) we infer that $\mu_r < 0$, and that there is a solution provided $|\mu_i| < 9H^2\omega\sqrt{\rho^2-\omega^2}/(4R\rho)$. In this case the current grows exponentially,

$$J_\phi = J_0^{(2)}\mathrm{e}^{(-3+\alpha_++\alpha_-)Ht}\,, \qquad (23)$$

with

$$J_0^{(2)} = (\phi_{+1}^{(0)}\phi_{-2}^{(0)} - \phi_{+2}^{(0)}\phi_{-1}^{(0)})H(\alpha_+-\alpha_-)\,. \qquad (24)$$

A typical evolution of the current (23–24) is shown in figure 2.

By making use of Eq. (16) we can relate the $\phi_{+1}^{(0)}$ and $\phi_{+2}^{(0)}$ such that the prefactor in Eqs. (21) and (24) becomes,

$$\phi_{+1}^{(0)}\phi_{-2}^{(0)} - \phi_{+2}^{(0)}\phi_{-1}^{(0)} = -\frac{4\phi_{+1}^{(0)}\phi_{+2}^{(0)}}{\mu_i}\sqrt{\frac{\rho^2}{\rho^2-\omega^2}\mu_r^2+\mu_i^2}\,. \qquad (25)$$

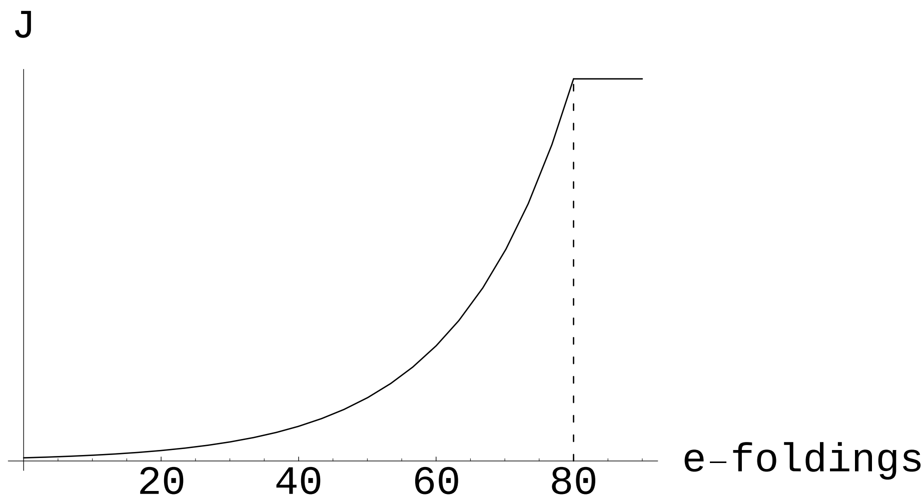

FIG. 2: The scalar charge density during inflation. Unlike in figure 1 no oscillations are present and the charge density grows exponentially.

Note that this relation and the above analysis are valid only in the presence of CP violation when $\mu_i \neq 0$. Indeed, Eq. (13) implies that for $\mu_i = 0$ the evolution of ϕ_+ and ϕ_- decouple and $\phi_+^{(0)}$ and $\phi_-^{(0)}$ can be specified independently. In this case there is no CP violation in the lagrangian. Nevertheless, even in this case an initial charge can be amplified provided there is a growing solution for the current. The solution is however always oscillatory and oscillations are rapid. On the other hand, amplification within a CP violating theory is more efficient and – as we have seen above – for a certain choice of parameters amplification is purely exponential.

In the above analysis we have assumed that R is varying very slowly and expansion is close to exponential. Most inflationary models, including inflation driven by a Brans-Dicke scalar field, are nearly exponential (powerlaw). Our preliminary investigation of powerlaw inflation indications that the current enhancement persists, albeit in that case it is weaker.

An equivalent charge amplification can be obtained during inflation with the replacement of $\mu^2 R \to m^2$, with m being the scalar mass. In that case CP violation remains operative after inflation, resulting in nontrivial effects in radiation and matter epochs. On the other hand, the curvature coupling is turned off in postinflationary radiation era, which makes the mechanism presented here in some sense simpler. Moreover, in the massive scalar case, no clear connection to the gravitational sector exists. Thus unless one finds another motivation for the scalar field (like e.g. flat directions in supersymmetric affleck-Dine baryogenesis models [2, 3]), the main impetus for this scalar-field baryogenesis mechanism would be lost.

After inflation the Universe enters a preheating epoch. The charged current can decay through a tree level coupling to ordinary matter in a charge conserving manner. Since the Brans-Dicke scalar is not charged under the gauge group of the Standard Model, it cannot couple directly to charged fermions and leptons of the Standard Model. It can however couple to neutrinos [8] *via* the Yukawa-type lagrangian,

$$\mathcal{L}_{\rm int} \supset -\sum_{i,j=1}^{3}(\bar{\nu}_R^i y_{ij}\phi\nu_L^j + {\rm h.c.})\,. \tag{26}$$

Note that when ϕ acquires an expectation value, this lagrangian generates Dirac neutrino masses. The interaction (26) induces the following tree level decay processes

$$\phi \to \bar{\nu}_L^i + \nu_R^j\,, \qquad \phi^* \to \bar{\nu}_R^i + \nu_L^j\,, \tag{27}$$

with the coupling strengths y_{ij} and y_{ij}^*, respectively, such that upon decay the scalar charge is converted into a lepton charge of a similar amplitude

$$J_\phi = Q \simeq L_0\,. \tag{28}$$

When sphalerons become active (below about $T \sim 10^{12}$ GeV, which corresponds to the sphaleron equilibration temperature), the lepton charge (28) gets converted into a net baryon (and lepton) number according to Eq. (10), with $(B-L)_0 \equiv -L_0$. For example, for $n_f = 3$ and $n_H = 1$, Eq. (10) gives $B = -(8/37)L_0 \simeq 0.2Q$.

During preheating a lot of entropy is produced, which results in an entropy density of the order $s \sim T^3$. Assuming instant reheating, the reheat temperature is $T \sim \sqrt{HM_P} \sim V^{1/4}$, such that the entropy density is at most, $s \sim V^{3/4}$. From this and the observed baryon-to-entropy ratio [13, 14],

$$\frac{n_B}{s} = (8.7 \pm 0.3) \times 10^{-11}\,, \tag{29}$$

we conclude that the required scalar current at the end of inflation is of the order

$$J_\phi \sim 10^{-9}V^{3/4}\,, \tag{30}$$

where $V \sim (10^{16}\ {\rm GeV})^4$ is the potential energy driving inflation. A comparison with the entropy density at the end of inflation, which is of the order $s_H \sim H^3$, is also instructive. The entropy production during preheating is of the order, $s_T \sim V^{3/4} \sim (V^{1/4}/H)^3 s_H \sim (M_P/H)^{3/2} s_H$, which is in typical one field scalar inflationary models of the order $10^8 s_H$. In these models we therefore need to produce a scalar charge density at the end of inflation of the order,

$$J_\phi|_{\rm end\ infl} \sim 10^{-1}H^3\,. \tag{31}$$

When compared with the above estimate of the expected magnitude of scalar charge fluctuations in the primordial preinflationary chaos, the required postinflationary current (31) seems easily feasible within the mechanism presented here.

One of the merits of the model is that it fits well into a model of extended inflation. Equation (7) resembles the lagrangian of curvature coupled extended inflation as proposed by Laycock and Liddle [15], except for the addition of the mixing term due to CP violation and the absence of an inflaton potential. Depending on the sign of μ_r, the fields $\phi_\pm$ assume the role of the inflaton or of the BD scalar. An example of the inflationary potential in which the inflaton couples to the Ricci scalar is,

$$V(\phi) = \lambda(|\phi|^2 - |\phi_0|)^2 = \sum_\pm \lambda(\phi_\pm^2 - |\phi_0|^2)^2 + 2\lambda\phi_+^2\phi_-^2 \tag{32}$$

where λ denotes a dimensionless quartic coupling. Since the valley of the potential (32) corresponds to the inflaton, for a negative μ_r, ϕ_- corresponds to the inflaton while ϕ_+ is a BD scalar. As we argued above, this choice for μ_r also provides a large parameter space for charge amplification. The BD scalar acquires a potential, which is however steeper than that of the inflaton. This may provide a way to anchoring the BD scalar to a fixed value today thus helping to make the theory compatible with experiments.

IV. CONCLUSION

By generalizing Brans-Dicke (BD) theory to a theory with a complex scalar field (2), we have constructed a baryogenesis model operative during inflation. In its most general disguise our BD theory contains one physical CP violating phase which can be used to amplify the corresponding scalar current during inflation.

Our starting assumption is that at the beginning of inflation the Universe contains domains of several Hubble volumes which contain a small amount of initial scalar charge. During inflation these domains are stretched to a volume larger than our visible universe, and at the same time the charge density gets (exponentially) enhanced. When inflation ends, CP violation turnes off and the Universe heats up, increasing the entropy of the Universe by a large factor. The scalar charge decays in a charge conserving way to ordinary matter through a tree level coupling, creating a net lepton and/or baryon number. The lepton minus baryon number is then processed via sphalerons to the lepton and baryon number of a magnitude which we observe today.

We have also shown that our mechanism can be incorporated into an extended inflationary model with a minimal adaptation when considering the real and imaginary part of the field as two independent real fields. The real part acts like a Brans-Dicke scalar while the imaginary part is the inflaton.

The theory can provide a natural potential for the Brans-Dicke scalar, potentially anchoring the expectation value today and loosening the constraints on the theory. More detailed calculations are needed to further investigate these proposals and compare them with the experimental bounds.

The main attractive feature of the generalised Brans-Dicke theory we consider here is in the fact that a model for baryogenesis is implemented in the theory of gravity itself. A related model, dubbed gravitational baryogenesis, was considered in Ref. [16], where baryons are produced via a dimension six operator which couples a standard model fermionic current to the gradient of the Ricci scalar. The authors argue that in homogeneous limit this interaction can be thought of as a chemical potential for baryon number of the form, $\mu_B \propto \dot{R} \propto \rho^{3/2}$, which therefore decays as $1/a^6$ in radiation era and as $1/a^{9/2}$ in matter era. That means that, unless the Universe starts at very high energy scales in a very asymmetric state, that mechanism does not work. Some may consider that as a feature which is not very desirable. Our mechanism does not suffer from such a problem since we require only a tiny preinflationary charge density which extends over several Hubble volumes, and which may as well be generated by statistical fluctuations in preinflationary primordial chaos.

Another interesting aspect of our model is that baryogenesis takes place during inflation, a very important era in the evolution of the universe not much considered in combination with baryogenesis. The CP and charge violation in the theory is turned off when the universe enters radiation era, such that after inflation the produced charge freezes in. It would be of interest to further investigate the model in a more general Scalar-Tensor theory of gravity or explore the analog in string theory with the dilaton.

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